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Synchrotron-Betatron Parametric Instability in Free-Electron Lasers

B. HAFIZI* AND A. TING

Beam Physics Branch
Plasma Physics Division

**Science Applications International Corp.*
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Analysis of the synchrotron-betatron coupling in a free-electron laser indicates that the particle motion is susceptible to an instability which is of a fundamentally different physical character to those examined previously. The domain of this instability, which is of a parametric type and causes an exponential growth of the betatron amplitude, is identified and the growth rate therein determined.

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CONTENTS

SYNCHROTRON-BETATRON PARAMETRIC INSTABILITY IN FREE-ELECTRON LASERS

I. Introduction

All electron beams have finite emittance; that is, there is a spread in the particle velocities transverse to the primary direction of motion. On traversing a wiggler, the inevitable presence of transverse gradients in the magnetic field lead to betatron oscillations of the particles. On the other hand, the beating of the wiggler and optical fields forms the ponderomotive potential wherein the particles perform synchrotron oscillations. Coupling of these two basic oscillatory degrees of freedom in a free-electron laser (FEL) may be deleterious to its efficient operation due either to detrapping out of the ponderomotive potential well or growth of betatron oscillation amplitude and emittance.

In Ref. 1 it is shown that the curvature of the optical wavefronts couples the synchrotron and betatron oscillations of the electrons and that under the proper resonance condition the amplitude of the synchrotron motion increases unstably. In subsequent work, the authors of Ref. 2 showed that a similar betatron-induced forcing of the synchrotron motion results in the case of planar optical wavefronts provided the wiggler field is tapered.

In what follows it is shown that the particle motion is additionally susceptible to an instability which is dependent neither on the curvature of the wavefronts nor on a tapered wiggler field. The instability arises from the energy-dependence of the betatron wavenumber. As a result when the electrons undergo synchrotron oscillations the betatron wavenumber is modulated at the synchrotron period, resulting in a parametric instability. This leads to an exponential growth of the betatron oscillation amplitude

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and velocity. The region of parameter space wherein this instability is operative is identified and an analytical expression for the growth rate is obtained.

II. Wiggle-average equations of motion

It is assumed that the wiggler is linearly polarized, with the vector potential given by

$$A_w = A_w \cosh(k_w y) \sin(k_w z)\, e_x,$$

where A_w is constant, $2\pi/k_w$ is the wiggler period and e_x is the unit vector along the x-axis. The vector potential of the linearly-polarized optical field is taken to correspond to the fundamental vacuum Gaussian mode:

$$A_s = A_s \cos(\frac{\omega}{c} z - \omega t + \phi)\, e_x,$$

where ω is the angular frequency, c is the speed of light in vacuo,

$$\phi(y,z) = \omega y^2/(2cR) - \frac{1}{2}\arctan(z/z_R) + \phi_o, \tag{1a}$$

$$R = z\left(1 + (z_R/z)^2\right) \tag{1b}$$

is the radius of curvature of the wavefronts,

$$z_R = \frac{\omega w_o^2}{2c} \tag{1c}$$

is the Rayleigh range, ϕ_o is a constant,

$$A_s = \hat{A}_s \left(\frac{w_o}{w(z)}\right)^{1/2} \exp\left(-y^2/w^2(z)\right), \tag{1d}$$

$\hat{A}_s$ is a constant, and $w(z)$, the spot size at z, is related to the waist w_o at $z = 0$ by

$$w(z) = w_o\left[1 + (z/z_R)^2\right]^{1/2}. \tag{1e}$$

It should be remarked that the amplitude and phase of the optical field are assumed to be constant in time.

Following the standard procedure,[3] the equations of motion in terms of the independent variable z are deriveable from the Hamiltonian function $-p_z(y,p_y;t,-E;z)$, which is given by

$$p_z = \frac{E}{c} - \frac{m^2c^3}{2E}\left(1 + (p_y/mc)^2 + a_x^2\right), \tag{2}$$

where E, the total energy of an electron of rest-mass m and charge $-|e|$, is canonically conjugate to the time t, p_y is the momentum conjugate to the coordinate y, and a_x is the normalized, total vector potential:

$$a_x = a_w\cosh(k_w y)\sin(k_w z) + a_s\cos(\frac{\omega}{c} z - \omega t + \phi), \tag{3}$$

with $a_{w,s} = |e|A_{w,s}/(mc^2)$. Since the vector potential is not an explicit function of the coordinate x, the associated canonical momentum P_x is a constant. Equation (2) has been derived by choosing $P_x = 0$, and omitting terms on the order of $(mc^2/E)^{-2}$ and higher in the expansion of the Hamiltonian.

Upon squaring the expression in Eq. (3), averaging over the wiggler period and retaining only the slowly-varying contribution to the ponderomotive potential, it is found that

$$p_z = \frac{E}{c} - \frac{m^2c^3}{2E}\Big\{1 + (p_y/mc)^2 + \frac{1}{2}\, a_w^2\cosh^2(k_w y) + a_w a_s \cosh(k_w y)$$

$$\times \sin\left[(\frac{\omega}{c} + k_w)z - \omega t + \phi\right]\Big\}.$$

Hamilton's equations are then given by

$$\frac{dy}{dz} = v, \tag{4a}$$

$$\frac{dv}{dz} = -k_\beta^2\left[\frac{\sinh(2k_w y)}{2k_w}\right], \tag{4b}$$

$$\frac{d\psi}{dz} = k_w - \frac{\omega}{2\gamma^2 c}\left[1 + \frac{1}{2}a_w^2\cosh^2(k_w y) + (\gamma v)^2\right], \tag{4c}$$

$$\frac{d\gamma}{dz} = \frac{-\omega}{2\gamma c}a_w a_s f_B \cosh(k_w y)\cos(\psi + \phi), \tag{4d}$$

where $\gamma = E/(mc^2)$, $\psi = (\omega/c + k_w)z - \omega t$, $v = p_y/(\gamma mc)$, $f_B = J_o(\xi) - J_1(\xi)$ is the usual difference of Bessel functions,[4] $\xi = (a_w/2)^2/(1 + a_w^2/2)$, and

$$k_\beta = \frac{a_w k_w}{\sqrt{2}\,\gamma} \tag{5}$$

is the betatron wavenumber. In arriving at Eq. (4b) it has been assumed that

$$\max\left\{\frac{a_s}{a_w}, \frac{\gamma a_s}{1 + a_w^2/2}, \frac{a_s}{a_w(k_w w_o)^2}\right\} << 1, \tag{6}$$

where w_o is the waist of the optical field, Eq. (1e). Additionally, the wave term, proportional to a_s/γ^2, has been dropped from Eq. (4c).

III. Analysis of synchrotron-betatron coupling

Variation of the optical phase ϕ across the wavefronts underlies the instability investigated in Ref. 1. Specifically, the first term in (1a) leads to a linear increase with z of the synchrotron oscillation amplitude of deeply trapped electrons. Assuming the radius of curvature R of the wavefronts to be constant, the growth rate of this instability is proportional to y_β^2/R, where y_β is the amplitude of the betatron

oscillation. In what follows it is assumed that $y_\beta^2/R \to 0$. This is the case for electrons confined to the central portion of the electron beam, or in the region where the curvature of the wavefronts is small. Assuming further that the z-variation of ϕ is small on the scale length of the instability to be discussed, for a synchronous electron Eqs. (4c) and (4d) imply

$$\gamma = \hat{\gamma} = \text{constant}, \tag{7a}$$

$$\hat{\psi} + \phi = (n + \tfrac{1}{2})\pi, \qquad n = 0, \pm 1, \pm 2, \ldots \tag{7b}$$

whence, for $|k_w y| \ll 1$, Eqs. (4a) and (4b) integrate to

$$y = y_\beta \cos(\hat{k}_\beta z) + \hat{k}_\beta^{-1} v_\beta \sin(\hat{k}_\beta z), \tag{8}$$

where, from (5), $\hat{k}_\beta = a_w k_w/\sqrt{2}\hat{\gamma}$ is the betatron wavenumber for a synchronous electron, and y_β, v_β are constants. Upon substituting (8) into Eq. (4c) one obtains

$$\frac{1}{2} a_w^2 \cosh^2(k_w y) + (\gamma v)^2 \simeq \frac{1}{2} a_w^2 \left[1 + (k_w y)^2\right] + (\gamma v)^2$$
$$\to \frac{1}{2} a_w^2 + \hat{\gamma}^2 \left(\hat{k}_\beta^2 y_\beta^2 + v_\beta^2\right),$$

which is a constant in z, whence the synchronous energy follows from Eq. (4c) by setting $d\hat{\psi}/dz = 0$:

$$\hat{\gamma}^2 = \frac{1 + \frac{1}{2} a_w^2 \left[1 + (k_w y_\beta)^2\right]}{2ck_w/\omega - v_\beta^2}. \tag{9}$$

For an electron in the vicinity of the synchronous particle, $\gamma = \hat{\gamma} + \delta\gamma$, $\psi = \hat{\psi} + \delta\psi$, and Eqs. (4c) and (4d) imply synchrotron oscillations of $\delta\gamma$ and of $\delta\psi$. Through the γ-dependence of the betatron

wavenumber, Eq. (5), it follows from Eqs. (4a) and (4b) that under an appropriate resonance condition a parametric instability of the coupled betatron and synchrotron oscillations is then possible.

To examine this, the expression in (8) is generalized to

$$y = y_\beta(z)\cos\left[(\hat{k}_\beta + \varepsilon)z\right] + (\hat{k}_\beta + \varepsilon)^{-1} v_\beta(z)\sin\left[(\hat{k}_\beta + \varepsilon)z\right], \tag{10}$$

where y_β and v_β are now functions of z and ε is a small wavenumber shift to be determined. Assuming $|dy_\beta/dz| \sim \varepsilon y_\beta$, $|dv_\beta/dz| \sim \varepsilon v_\beta$, and omitting terms on the order of ε^2, substitution of (10) into Eqs. (4a) and (4b) yields

$$\begin{aligned}(dv_\beta/dz - \varepsilon\hat{k}_\beta y_\beta)\cos\left[(\hat{k}_\beta + \varepsilon)z\right] - (\hat{k}_\beta dy_\beta/dz + \varepsilon v_\beta)\sin\left[(\hat{k}_\beta + \varepsilon)z\right] \\ = \hat{k}_\beta^2 \left\{y_\beta\cos\left[(\hat{k}_\beta + \varepsilon)z\right] + (\hat{k}_\beta + \varepsilon)^{-1} v_\beta \sin\left[(\hat{k}_\beta + \varepsilon)z\right]\right\}(\delta\gamma/\hat{\gamma}).\end{aligned} \tag{11}$$

To obtain an expression for $\delta\gamma(z)$ to be used in Eq. (11), Eqs. (4c) and (4d) are perturbed about the synchronous values (7a) and (7b):

$$\begin{aligned}\frac{d}{dz}\,\delta\psi = k_w - \frac{\omega}{2c\hat{\gamma}^2}\left[1 + \frac{1}{2}\,a_w^2\left(1 + k_w^2 y^2\right) + \hat{\gamma}^2 v^2\right] \\ + \frac{\omega}{c\hat{\gamma}^3}\left[1 + \frac{1}{2}\,a_w^2\left(1 + k_w^2 y^2\right)\right]\delta\gamma,\end{aligned} \tag{12a}$$

$$\frac{d}{dz}\,\delta\gamma = (-1)^n\,\frac{\omega a_w a_s f_B}{2\hat{\gamma}c}\left(1 + \frac{1}{2}\,k_w^2 y^2\right)\delta\psi, \tag{12b}$$

where n denotes an integer from the set defined in (7b). Substituting (10) and the forms

$$\delta\gamma = \Delta\gamma(z)\cos\left[2(\hat{k}_\beta + \varepsilon)z + \eta\right], \tag{13a}$$

$$\delta\psi = \Delta\psi(z)\sin\left[2(\hat{k}_\beta + \varepsilon)z + \eta\right], \tag{13b}$$

where η is a constant, into (12) and assuming that

$$k_w y_\beta << \hat{\gamma}^{-1/2}, \tag{14}$$

one obtains an identity whose consistency requires that $\Delta\gamma$ and $\Delta\psi$ be independent of z and that

$$2(\hat{k}_\beta + \varepsilon) = \hat{k}_{syn}, \tag{15}$$

where

$$\hat{k}_{syn} = \left[\frac{\omega^2 a_w a_s f_B}{2\hat{\gamma}^4 c^2}\left(1 + \frac{1}{2} a_w^2\right)\right]^{1/2}$$

is the synchrotron wavenumber. In addition, it is required that the integer n in Eq. (12b) be odd. It should be noted that the constraint indicated by (14) is consistent with the neglect of the radiation field term in Eq. (4c) and the parameter restriction in (6).

Now substituting (13a) into Eq. (11) one again obtains an identity. Assuming y_β, $v_\beta \sim \exp(\lambda z)$, the consistency condition reduces to

$$\lambda = \pm \left\{\left[\frac{1}{2} \hat{k}_\beta (\Delta\gamma/\hat{\gamma})\cos\eta\right]^2 - \varepsilon^2\right\}^{1/2}, \tag{16}$$

which is the growth rate of the instability.

IV. Discussion

It is clear that the source of the instability discussed here lies in the energy dependence of the betatron wavenumber, Eq. (5). Previous studies of the synchrotron-betatron coupling have neglected this dependence. The results of Refs. 1 and 2 indicate a linear growth in the amplitude of the synchrotron motion ($\Delta\gamma$, $\Delta\psi$) for deeply trapped electrons.

On the other hand, for the parametric instability discussed herein, $\Delta\gamma$ and $\Delta\psi$ are constants, whereas y_β and v_β grow exponentially fast. In all three cases, however, the resonance condition is essentially $\hat{k}_{syn} \simeq 2\hat{k}_\beta$.

From Eqs. (15) and (16) it follows that for a sufficiently large wavenumber mismatch ε the growth rate vanishes. Maximum growth is obtained for $\varepsilon = 0$:

$$\lambda_{max} = \frac{1}{2}\,\hat{k}_\beta\left|(\Delta\gamma/\hat{\gamma})\cos\eta\right|. \tag{17}$$

Since, at the order of approximation employed here, $\Delta\gamma$ and $\Delta\psi$ are constants, the parametric instability is expected to affect the extraction efficiency of the FEL mainly through increased electron beam radius and particle detrapping from the ponderomotive potential, such as that described in Refs. 1 and 2, would result once the betatron amplitude is sufficiently large. It must be remarked, however, that if the variation of the optical field amplitude or phase over the scale length indicated by (17) is large enough, then the resonance (15) may be crossed without an appreciable effect on the efficiency.[5,6]

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University of New Mexico
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Dr. Philip Morton
Stanford Linear Accelerator Center
P.O. Box 4349
Stanford, CA 94305

Prof. J. Nation
224 Phillips Hall
School of Elec. Eng.
Cornell University
Ithaca, NY 14850

Dr. George Neil
TRW
One Space Park
Redondo Beach, CA 90278

Dr. Kelvin Neil
Lawrence Livermore National Lab.
Code L-321, P.O. Box 808
Livermore, CA 94550

Dr. Brian Newnam
MSJ 564
Los Alamos National Scientific Lab.
P.O. Box 1663
Los Alamos, NM 87545

Dr. T. Orzechowski
L-436
Lawrence Livermore National Lab.
P. O. Box 808
Livermore, CA 94550

Prof. E. Ott
Department of Physics
University of Maryland
College Park, MD 20742

OUSDRE (R&AT)
Room 3D1067, The Pentagon
Washington, D.C. 20301

Dr. Robert B. Palmer
Brookhaven National Laboratories
Associated Universities, Inc.
Upton, L.I., NY 11973

Dr. J. Palmer
Hughes Research Laboratory
Malibu, CA 90265

Dr. Richard H. Pantell
Stanford University
Stanford, CA 94305

Dr. Dennis Papadopoulos
Astronomy Department
University of Maryland
College Park, Md. 20742

Dr. John A. Pasour
Mission Research Laboratory
8560 Cinderbed Road
Suite 700
Newington, VA 22122

Dr. C. K. N. Patel
Bell Laboratories
Murray Hill, NJ 07974

Dr. Claudio Pellegrini
Brookhaven National Laboratory
Associated Universities, Inc.
Upton, L.I., NY 11973

Dr. S. Penner
Center for Radiation Research
Natl. Inst. of Standards and Tech.
Gaithersburg, MD 20899

Dr. M. Piestrup
Adelphi Technology
13800 Skyline Blvd. No. 2
Woodside, CA 94062

Dr. D. J. Pistoresi
Boeing Aerospace Company
P. O. Box 3999
Seattle, WA 98124-2499

Major E. W. Pogue
SDIO
The Pentagon, T-DE Rm. 1E180
Washington, DC 20301-7100

Major Donald Ponikvar
U. S. Army SDC
P. O. Box 15280
Arlington, VA 22245-0280

Dr. Donald Prosnitz
Lawrence Livermore National Lab.
Box 5511 L-626
Livermore, CA 94550

Dr. D. C. Quimby
Spectra Technology
2755 Northup Way
Bellevue, WA 98004

Dr. G. Ramian
Quantum Institute
University of California
Santa Barbara, CA 93106

Dr. M. Reiser
University of Maryland
Department of Physics
College Park, MD 20742

Dr. S. Ride
Arms Control
Stanford University
Stanford, CA 94305

Dr. C. W. Roberson
Office of Naval Research
Code 112S
800 N. Quincy Street
Arlington, VA 22217

Dr. K. Robinson
Spectra Technology
2755 Northup Way
Bellevue, WA 98004

Dr. Marshall N. Rosenbluth
Dept. of Physics
B-019
Univ. of Calif., San Diego
LaJolla, CA 92093

Dr. J. B. Rosenzweig
The Inst. for Accelerator Physics
Department of Physics
University of Wisconsin-Madison
Madison, WI 53706

Dr. N. Rostoker
Department of Physics
University of California at Irvine
Irvine, CA 92717

Dr. A. Saxman
Los Alamos National Scientific Lab.
P. O. Box 1663, MSE523
Los Alamos, NM 87545

Dr. E. T. Scharlemann
L626
Lawrence Livermore National Laboratory
P. O. Box 808
Livermore, CA 94550

Prof. S. P. Schlesinger
Dept. of Electrical Engineering
Columbia University
New York, NY 10027

Dr. Howard Schlossberg
AFOSR
Bolling AFB
Washington, D.C. 20332

Dr. George Schmidt
Stevens Institute of Technology
Physics Department
Hoboken, NJ 07030

Dr. M. J. Schmitt
Los Alamos National Laboratory
P. O. Box 1663
Los Alamos, NM 87545

Dr. H. Schwettmann
Phys. Dept. & High Energy
Physics Laboratory
Stanford University
Stanford, CA 94305

Dr. Marlan O. Scully
Dept. of Physics & Astronomy
Univ. of New Mexico
800 Yale Blvd. NE
Albuquerque, NM 87131

Dr. S. B. Segall
KMS Fusion
3941 Research Park Dr.
P.O. Box 1567
Ann Arbor, MI 48106

Prof. P. Serafim
Northeastern University
Boston, MA 02115

Dr. A. M. Sessler
Lawrence Berkeley Laboratory
University of California
1 Cyclotron Road
Berkeley, CA 94720

Dr. W. Sharp
L-626
Lawrence Livermore National Laboratory
P. O. Box 808
Livermore, CA 94550

Dr. Earl D. Shaw
Bell Laboratories
600 Mountain Avenue
Murray Hill, NJ 07974

Dr. R. L. Sheffield
Los Alamos National Laboratory
P.O. Box 1663
Los Alamos, NM 87545

Dr. D. Shoffstall
Boeing Aerospace Company
P.O. Box 3999
Seattle, WA 98124

Dr. Jack Slater
Spectra Technology
2755 Northup Way
Bellevue, WA 98004

Dr. Todd Smith
Hansen Labs
Stanford University
Stanford, CA 94305

Dr. R. Sudan
Lab. of Plasma Studies
Cornell University
Ithaca, NY 14850

Dr. David F. Sutter
ER 224, GTN
Department of Energy
Washington, D.C. 20545

Dr. T. Tajima
Institute for Fusion Studies
University of Texas at Austin
Austin, TX 78712

Dr. R. Temkin
Mass. Institute of Technology
Plasma Fusion Center
Cambridge, MA 02139

Dr. L. Thode
Los Alamos National Laboratory
P. O. Box 1663
Los Alamos, NM 87545

Dr. Norman H. Tolk
Physics Department
Vanderbilt University
Nashville, TN 37240

Dr. Kang Tsang
Science Applications Intl. Corp.
1710 Goodridge Dr.
McLean, VA 22102

Dr. H. S. Uhm
Naval Surface Warfare Center
White Oak Lab.
Silver Spring, MD 20903-5000

Under Secretary of Defense (R&D)
Office of the Secretary of Defense
Room 3E1006, The Pentagon
Washington, D.C. 20301

Dr. John E. Walsh
Wilder Laboratory
Department of Physics (HB 6127)
Dartmouth College
Hanover NH 03755

Dr. Jiunn-Ming Wang
Brookhaven National Laboratories
Associated Universities, Inc.
Upton, L.I., NY 11973

Dr. Roger W. Warren
Los Alamos National Scientific Lab.
P.O. Box 1663
Los Alamos, NM 87545

Dr. J. Watson
Los Alamos National Laboratory
P. O. Box 1663
Los Alamos, NM 87545

Dr. Mark Wilson
Natl. Inst. of Standards and Tech.
Bldg. 245, Rm. B-119
Gaithersburg, MD 20899

Dr. J. Wurtele
M.I.T.
NW 16-234
Plasma Fusion Center
Cambridge, MA 02139

Dr. Ming Xie
Dept. of Physics
Stanford University
Stanford, CA 94305

Dr. Yi-Ton Yan
Stanford Linear Accelerator Center
P. O. 4349
Stanford, CA 94305

Dr. Simon S. Yu
Lawrence Livermore National Laboratory
P. O. Box 808
Livermore, CA 94550

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Naval Research Laboratory
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